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Raman-induced spectral shifts in optical fibers: general theory based on the moment method

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Abstract

The Raman-induced spectral shift of ultrashort pulses, having its origin in the phenomenon of intrapulse Raman scattering occurring inside silica fibers, has been mostly studied in the context of optical solitons. We use the moment method to show that such a spectral shift occurs both in the normal and anomalous-dispersion regimes and depends not only on the pulse width but also on the frequency chirp associated with an optical pulse. The magnitude of spectral shift depends on the history of pulse width changes and saturates to a constant value if the pulse broadens considerably during propagation.

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1. Introduction

It was discovered in 1985 numerically [1] that the spectrum of an ultrashort optical pulse can shift toward longer wavelengths (a "red" shift) when the pulse propagates in the anomalous-dispersion regime of an optical fiber as an optical soliton. Such a spectral shift was observed in a 1986 experiment [2] by using a stabilized, mode-locked laser capable of emitting pulses shorter than 1 ps. It was called the soliton self-frequency shift because pulses whose spectrum was red-shifted were propagating as solitons inside the optical fiber used in the experiment. In fact, Gordon used a perturbation theory of solitons for predicting the magnitude of the spectral shift and its dependence on the pulse and fiber parameters [3]. Physically, the spectral shift is attributed to intrapulse Raman scattering (IRS), a phenomenon in which high-frequency components of an optical pulse pump the low-frequency components of the same pulse, thereby transferring energy to the red side through stimulated Raman scattering. Since 1986, the Raman-induced frequency shift (RIFS) has been studied extensively for both constant dispersion and dispersion-managed fibers but mostly in the context of solitons [4–9].

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From a physical standpoint, it is hard to see why the Raman-induced red shift of ultrashort pulses should require the formation of solitons. The IRS phenomenon should occur for any optical pulse, irrespective of whether it propagates in the normal- or anomalous-dispersion regime of an optical fiber. It should also be affected by the frequency chirp, if the input pulse is chirped. In this paper we develop a general theory of IRS using the moment method and use it to study the effect of frequency chirp on the RIFS in the cases of both normal and anomalous dispersion. In Section 2 we discuss the moment method and show how it can be used to calculate the RIFS. The case of sech-shaped pulses is discussed in Section 3 where we show that our theory reduces to that of Gordon [3] when the pulse propagates as a standard soliton only if the RIFS is relatively small so that it does not affect the soliton itself. In Section 4 we consider input pulses in the form of a chirped Gaussian pulse and discuss the effects of frequency chirp on RIFS. The main results are summarized in Section 5.

2. The moment method

Propagation of ultrashort optical pulses in optical fibers is governed by a generalized nonlinear Schrödinger (NLS) equation of the form [9]

$$\frac{\partial A}{\partial z}+\frac{\alpha}{2}A+\frac{\mathrm{i}\beta_2}{2}\frac{\partial^2 A}{\partial t^2}-\frac{\beta_3}{6}\frac{\partial^3 A}{\partial t^3} = \mathrm{i}\gamma\left(|A|^2A+\frac{\mathrm{i}}{\omega_0}\frac{\partial}{\partial t}(|A|^2A)-T_{\mathrm{R}}A\frac{\partial |A|^2}{\partial t}\right), \quad (1)$$

where A is the slowly varying amplitude of the pulse envelope, α accounts for fiber losses, β_2 is the group-velocity dispersion (GVD) coefficient, β_3 is the third-order dispersion (TOD) parameter, γ is the nonlinear parameter responsible for self-phase modulation, ω_0 is the center frequency of the pulse spectrum and the Raman parameter T_{R} accounts for IRS. Eq. (1) is difficult to be solved analytically in general and must be solved either numerically or by using an approximate analytic technique. For completeness, we have included all higher-order terms in Eq. (1). If the pulse width exceeds 1 ps, the shock term containing ω_0 and the TOD term containing β_3 can sometimes be dropped.

The variational method is often used [10] for solving Eq. (1). However, the presence of the IRS term in Eq. (1) makes the variational technique unsuitable because the Lagrangian density needed for it does not exist when $T_{\mathrm{R}} \neq 0$. The fundamental reason behind this is that the IRS terms does not conserve energy as each red-shifted photon has less energy than the original one. For this reason, we adopt the moment method that has been used recently for calculating timing jitter in lightwave systems [11–13]. The basic idea is to treat the optical pulse like a particle [14] whose energy E, position T, and the frequency shift Ω (from the original carrier frequency) are defined as:

$$E=\int_{-\infty}^{\infty}|A|^2\,\mathrm{d}t, \quad (2)$$

$$T=\frac{1}{E}\int_{-\infty}^{\infty}t|A|^2\,\mathrm{d}t, \quad (3)$$

$$\Omega=\frac{\mathrm{i}}{2E}\int_{-\infty}^{\infty}\left(A^*\frac{\partial A}{\partial t}-A\frac{\partial A^*}{\partial t}\right)\mathrm{d}t, \quad (4)$$

where A satisfies Eq. (1). The root-mean square (RMS) width of the pulse is then defined as

$$\sigma^2=\frac{1}{E}\int_{-\infty}^{\infty}(t-T)^2|A|^2\,\mathrm{d}t. \quad (5)$$

The actual pulse width is related to the RMS width by a constant factor that depends on the pulse shape. We introduce one more moment related to the chirp of the pulse by the same constant factor using the definition

$$\tilde{C}=\frac{\mathrm{i}}{2E}\int_{-\infty}^{\infty}(t-T)\left(A^*\frac{\partial A}{\partial t}-A\frac{\partial A^*}{\partial t}\right)\mathrm{d}t. \quad (6)$$

Using Eq. (1) in Eqs. (2)–(6) the five moments E, T, Ω, σ and $\tilde{C}$ are found to evolve along the fiber according to the following set of five ordinary differential equations:

$$\frac{\mathrm{d}E}{\mathrm{d}z}=-\alpha E, \quad (7)$$

$$\frac{\mathrm{d}T}{\mathrm{d}z}=\beta_2\Omega+\frac{\beta_3}{2E}\int_{-\infty}^{\infty}\left|\frac{\partial A}{\partial t}\right|^2\mathrm{d}t+\frac{3\gamma}{2\omega_0 E}\int_{-\infty}^{\infty}|A|^4\,\mathrm{d}t, \quad (8)$$

$$\frac{d\Omega}{dz} = -\frac{T_R\gamma}{E}\int_{-\infty}^{\infty}\left(\frac{\partial|A|^2}{\partial t}\right)^2 dt - \frac{i\gamma}{E\omega_0} \times \int_{-\infty}^{\infty}\left(A^*\frac{\partial A}{\partial t} - A\frac{\partial A^*}{\partial t}\right)\frac{\partial}{\partial t}|A|^2\,dt, \tag{9}$$

$$\frac{d\sigma}{dz} = \frac{\beta_2\tilde{C}}{\sigma} + \frac{\beta_3}{2\sigma E}\int_{-\infty}^{\infty}(t-T)\left|\frac{\partial A}{\partial t}\right|^2 dt, \tag{10}$$

$$\begin{aligned}\frac{d\tilde{C}}{dz} = {} & \frac{\beta_2}{E}\int_{-\infty}^{\infty}\left|\frac{\partial A}{\partial t}\right|^2 dt \\ & + \frac{i\beta_3}{4E}\int_{-\infty}^{\infty}\left(\frac{\partial^2 A}{\partial t^2}\frac{\partial A^*}{\partial t} - \frac{\partial^2 A^*}{\partial t^2}\frac{\partial A}{\partial t}\right)dt \\ & + \frac{\gamma}{2E}\int_{-\infty}^{\infty}|A|^4\,dt - \frac{i\gamma}{E\omega_0}\int_{-\infty}^{\infty}(t-T)\frac{\partial|A|^2}{\partial t} \\ & \times\left(A^*\frac{\partial A}{\partial t} - A\frac{\partial A^*}{\partial t}\right)dt \\ & - \frac{i\gamma}{2E\omega_0}\int_{-\infty}^{\infty}|A|^2\left(A\frac{\partial A^*}{\partial t} - A^*\frac{\partial A}{\partial t}\right)dt \\ & - \frac{\gamma T_R}{E}\int_{-\infty}^{\infty}(t-T)\left(\frac{\partial|A|^2}{\partial t}\right)^2 dt. \end{aligned} \tag{11}$$

Eqs. (7)–(11) reduce the complexity of the problem but they are still not in a useful form because they depend on $A(z,t)$, which is not known until Eq. (1) is solved. If one has some knowledge of the pulse shape and its dependence on the five moments, the problem can be solved approximately. There are several situations in which pulse shape is known a priori with a good degree of approximation. For example, in the case of standard solitons, pulse shape can be assumed to maintain "sech" shape even if its width changes when the higher-orders terms are relatively small and can be treated as a small perturbation. As another example, pulse shape remains nearly Gaussian in a dispersion-managed fiber link under the same conditions [15]. In general, a Gaussian pulse can be assumed to maintain its shape during propagation inside optical fibers if the nonlinear length is much larger than the dispersion length [9] and higher-order effects are relatively weak. We consider these two cases in the following two sections.

3. Hyperbolic secant pulses

We first consider the propagation of standard solitons in a fiber with constant dispersion. Normally, standard solitons are unchirped in the absence of IRS. However, the chirp-free nature is not ensured when their spectrum shifts because of IRS. For this reason, we allow for a chirp on the input pulse but assume it to be small enough that the soliton shape does not change. In this case, the pulse amplitude can be written as

$$A(z,t) = \sqrt{\frac{E}{2\tau}}\,\mathrm{sech}\left(\frac{t-T}{\tau}\right)\exp[i\phi - i\Omega(t-T) - iC(t-T)^2/2\tau^2], \tag{12}$$

where the phase ϕ does not depend on t. The width parameter τ and the chirp parameter C appearing in this equation is related to the RMS width σ and the moment $\tilde{C}$, respectively, by a constant factor K such that $\tau^2 = K\sigma^2 = (12/\pi^2)\sigma^2$ and $C = K\tilde{C} = (12/\pi^2)\tilde{C}$. All pulse parameters represent local values and change along the fiber with z. We substitute Eq. (12) in Eqs. (7)–(11) and perform the integration over t. The five pulse parameters are then found to evolve as:

$$\frac{dE}{dz} = -\alpha E, \tag{13}$$

$$\frac{dT}{dz} = \beta_2\Omega + \frac{\beta_3}{2}\left[\Omega^2 + \left(1+\frac{\pi^2}{4}C^2\right)\frac{1}{3\tau^2}\right] + \frac{\gamma E}{2\omega_0\tau}, \tag{14}$$

$$\frac{d\Omega}{dz} = -\frac{4T_R\gamma E}{15\tau^3} + \frac{\gamma CE}{3\omega_0\tau^3}, \tag{15}$$

$$\frac{d\tau}{dz} = \frac{\beta_2 C}{\tau} + \beta_3\frac{C\Omega}{\tau}, \tag{16}$$

$$\frac{dC}{dz} = \left(\frac{4}{\pi^2}+C^2\right)\frac{\beta_2}{\tau^2} + \frac{12}{\pi^2}\beta_2\Omega^2 + \frac{2\gamma E}{\pi^2\tau} + \frac{\beta_3\Omega}{2\tau^2}\left(\frac{4}{\pi^2}+3C^2\right) + \frac{6}{\pi^2}\beta_3\Omega^3 + \frac{24\gamma\Omega E}{\pi^2\omega_0\tau}. \tag{17}$$

These equations look similar to those found when the variational method is used. In fact, they can also be obtained with the variational method when $T_R = 0$.

Consider first the special case of chirp-free solitons launched in a fiber whose losses are exactly compensated through distributed amplification such that losses vanish effectively ($\alpha = 0$). The pulse energy E then remains constant. If we ignore the higher-order effects except for IRS by setting $\omega_0 = 0$ and $\beta_3 = 0$ and use $C = 0$ in Eqs. (13)–(17), we find that τ remain constant along the fiber, as it should for solitons. Also, E and τ are not independent but related to each other by the soliton condition $L_D = L_{NL}$, where $L_D = \tau^2/|\beta_2|$ and $L_{NL} = (\gamma P_0)^{-1}$ are the dispersion and nonlinear length scales [9]. This condition can be obtained from Eq. (17) by setting $dC/dz = 0$ if we neglect the Ω term and relate the peak power P_0 of the solitons to the soliton energy using $E = 2P_0\tau$. Using the condition $L_D = L_{NL}$, we find that $E = 2|\beta_2|/(\gamma\tau)$. If we substitute this relation in Eq. (15), the RIFS evolves as

$$\Omega(z) = -\frac{8T_R|\beta_2|z}{15\tau^4}. \tag{18}$$

Eq. (18) is identical to the RIFS magnitude first estimated by Gordon using perturbation theory [3]. It shows that the RIFS increases linearly with distance but scales with pulse width as τ^{-4}, thereby becoming important only for pulses shorter than a few picoseconds. However, its derivation assumes that the soliton remains unchirped. From Eq. (17), C remains zero for solitons only if $\Omega = 0$. Eqs. (16) and (17) clearly show that C and τ both begin to change for standard solitons because of the RIFS. Thus, Eq. (18) is only valid in the limit in which the RIFS is small enough that it does not affect the soliton. We can find the validity condition for Eq. (18) in the absence of third-order dispersion and self-steepening, by requiring in Eq. (17) that $|\beta_2|\Omega^2 \ll 2\gamma E/(\pi^2\tau)$. Using $E = 2|\beta_2|/(\gamma\tau)$, this condition is equivalent to requiring $\Omega\tau \ll 1$. Noting that the spectral width of a pulse scales inversely with the pulse width τ, we conclude that Eq. (18) is valid as long as the RIFS remains a small fraction of the pulse spectral width. In many practical situations, RIFS becomes large enough that it exceeds the spectral width of pulse significantly.

We thus consider the more general case in which neither E nor τ remains constant along the fiber. The pulse energy E generally changes because of gain–loss variations introduced when losses are compensated periodically using optical amplifiers [15]. The soliton width τ begins to change as soon as the pulse becomes chirped ($C \neq 0$). Using $E = E_0 e^{-\alpha z}$, the total frequency shift is found by integrating Eq. (15) and is given by

$$\Omega(z) = -\frac{4T_R\gamma E_0}{15}\int_0^z \frac{e^{-\alpha z}}{\tau^3}\,dz + \frac{\gamma E_0}{3\omega_0}\int_0^z C(z)\frac{e^{-\alpha z}}{\tau^3}\,dz. \tag{19}$$

Note that the RIFS depends on the local pulse width as τ^{-3} and not as τ^{-4}, as suggested by Eq. (18). Of course, the z dependence of C and τ should be calculated by solving Eq. (16) which is coupled to the chirp Eq. (17), which in turn depends on Ω itself. It is this interdependence among τ, C and Ω that governs the eventual magnitude of the RIFS. The last term in Eq. (19) shows that the shock term also induces a spectral shift but its sign (red versus blue shift) and magnitude depend on how the chirp $C(z)$ changes along the fiber length.

As a numerical example, consider the propagation of solitons at wavelength $\lambda = 1.55$ μm with $\tau_0 = 50$ fs (full-width at half-maximum about 88 fs) in a 10-m-long, dispersion-shifted fiber with the GVD of 4 ps/km nm ($|\beta_2| = 5.1$ ps^2/km). Figs. 1 and 2 show the RIFS and pulsewidth τ as a function of distance z in the cases of anomalous and normal dispersion, respectively. The nonlinear parameter $\gamma = 1.994$ W^{-1} km^{-1} was calculated using an effective core area of 50 μm^2. Also $\alpha = 0.2$ dB/km and $\beta_3 = 0.1$ ps^3/km. Consider the case of anomalous dispersion first as it corresponds to the propagation of solitons. The solid curve in Fig. 1 shows the $C(0) = 0$ case that corresponds to standard solitons. The pulse width is indeed maintained in the beginning, as expected, but begins to increase after 2 m because of the RIFS and TOD effects. The magnitude of RIFS becomes comparable to the spectral width of the pulse (about 2 THz) at a distance of 2 m, and it begins to affect the soliton itself. Notice that Ω increases linearly initially up to a distance of 2 m but then begins to saturate as the

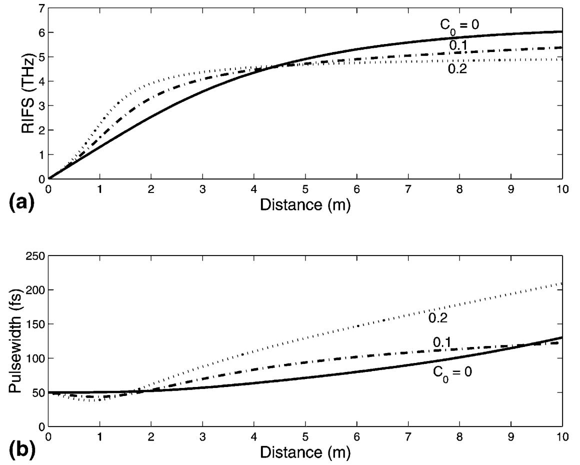

Fig. 1. Evolution of Raman-induced frequency shift (a) and pulse width (b) when sech-shaped pulses with $\lambda = 1.55$ μm and $T_0 = 50$ fs propagate inside a 10-m-long fiber exhibiting anomalous dispersion ($D = 4$ ps/km nm). The input chirp parameter C_0 varies in the range 0–0.2 for the three curves.

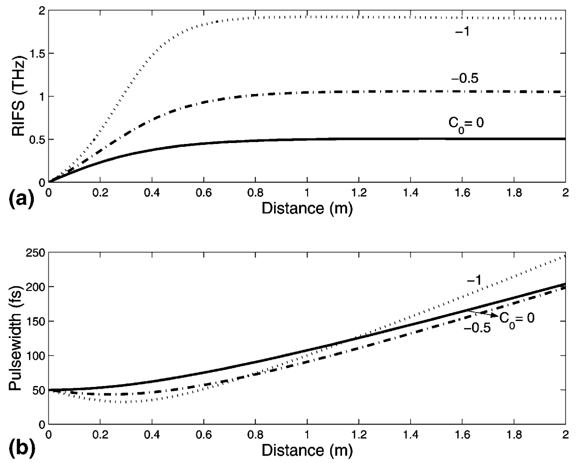

Fig. 2. Same as in Fig. 1 except that fiber exhibits normal dispersion ($D = -4$ ps/km nm). The input chirp parameter C_0 varies in the range −1 to 0 for the three curves.

pulse width increases. The use of Eq. (18) would be inappropriate under such conditions. The dashed and dash-dotted lines show that even a relatively small chirp affects the RIFS considerably. For positive values of C, the pulse is initially compressed, as expected for $\beta_2 C < 0$ [9], and then broadens after attaining its shortest width at a distance of about 1 m. For this reason Ω initially increases before saturating as the pulse broadens. The main point to note is that the chirp can increase the RIFS whenever $\beta_2 C < 0$, and fiber is not much longer than the length at which pulse becomes transform-limited. For $C < 0$, pulse begins to broaden immediately, and RIFS is reduced considerably.

Fig. 2 shows the RIFS and pulse width as a function of distance in the case of normal dispersion. The solid line again shows the case $C(0) = 0$. Since pulse begins to broaden right away, in contrast with the soliton case where pulse width remained constant for up to 2 m, Ω quickly saturates and is thus considerably smaller in magnitude in the case of normal GVD. The dashed and dash-dotted lines show that it can be enhanced by chirping the input Gaussian pulse such that $\beta_2 C_0 < 0$. The reason is easily understood by noting that the pulse can be compressed by a factor of $\sqrt{2}$ for $|C_0| = 1$, and the compression factor can be increased even more for large values of the chirp. As seen in Fig. 2, almost entire RIFS occurs within the first meter of the fiber, where pulse remains compressed and its magnitude is about three times larger for $|C_0| = 1$ compared with the for $C_0 = 0$ case. With sufficiently large chirp, the RIFS can even become comparable to that obtained in the case of anomalous dispersion. We thus conclude that RIFS can be made large enough to be measurable even in the case of normal GVD through proper chirp control.

4. Chirped Gaussian pulses

In this section we consider the case of a Gaussian pulse shape. In particular, we assume that a chirped Gaussian pulse is launched initially into the fiber and it maintains this shape during propagation even thought its width and chirp may change. The optical field can then be written as

$$A(z,t) = \sqrt{\frac{E}{\pi\tau}}\exp[-(1+\mathrm{i}C)(t-T)^2/2\tau^2 + \mathrm{i}\phi - \mathrm{i}\Omega(t-T)]. \tag{20}$$

The width parameter τ and the chirp parameter C appearing in this equation is again related to the

RMS width σ and the moment $\tilde{C}$ by a constant factor K such that $\tau^2 = K\sigma^2 = 2\sigma$ and $C = 2\tilde{C}$, respectively. Using Eq. (20) in Eqs. (7)–(11) and performing integrations, the five pulse parameters are found to evolve as:

$$\frac{\mathrm{d}E}{\mathrm{d}z} = -\alpha E, \tag{21}$$

$$\frac{\mathrm{d}T}{\mathrm{d}z} = \beta_2\Omega + \frac{\beta_3}{2}\left(\Omega^2 + \frac{1+C^2}{2\tau^2}\right) + \frac{3\gamma E}{\sqrt{8\pi}\omega_0\tau}, \tag{22}$$

$$\frac{\mathrm{d}\Omega}{\mathrm{d}z} = -\frac{T_\mathrm{R}\gamma E}{\sqrt{2\pi}\tau^3} + \frac{\gamma C E}{\sqrt{2\pi}\omega_0\tau^3}, \tag{23}$$

$$\frac{\mathrm{d}\tau}{\mathrm{d}z} = \frac{\beta_2 C}{\tau} + \frac{\beta_3 C\Omega}{\tau}, \tag{24}$$

$$\frac{\mathrm{d}C}{\mathrm{d}z} = \beta_2\left(\frac{1+C^2}{\tau^2}\right) + \frac{\gamma E}{\sqrt{2\pi}\tau} + 2\beta_2\Omega^2 + \beta_3\Omega^3 + \beta_3\left(\frac{1+C^2}{\tau^2}\right) + \frac{4\gamma\Omega E}{\sqrt{2\pi}\omega_0\tau}. \tag{25}$$

Following the method discussed in Section 3, the RIFS in the Gaussian case is given by

$$\Omega(z) = -\frac{T_\mathrm{R}\gamma E_0}{\sqrt{2\pi}}\int_0^z \frac{\mathrm{e}^{-\alpha z}}{\tau^3}\,\mathrm{d}z + \frac{\gamma E_0}{\sqrt{2\pi}\omega_0}\int_0^z C(z)\frac{\mathrm{e}^{-\alpha z}}{\tau^3}\,\mathrm{d}z, \tag{26}$$

where $C(z)$ and $\tau(z)$ should be found numerically by solving Eqs. (21)–(25). Eq. (26) should be compared with Eq. (19) found in the "sech" case. It is evident that the exact shape of the pulse has a relatively minor effect on the magnitude of the RIFS. In particular, the functional dependence on the local pulsewidth and local magnitude of loss remains exactly the same. Even the numerical factor of $(2\pi)^{-1/2} \approx 0.4$ in the Gaussian case is only slightly larger than the factor of $4/15 \approx 0.267$ found in the "sech" case. This feature indicates that even if the pulse shape deviates somewhat from the shape assumed in applying the moment method, our analysis should still provide a good estimate of the RIFS in practice.

Figs. 3 and 4 shows the RIFS and pulse width of Gaussian pulses as a function of distance in the cases of anomalous and normal dispersion, respectively, for the same 10-m long fiber and for the same set of parameters used for Figs. 1 and 2. For a fair comparison with the soliton case, the initial pulse energy is chosen to be $E_0\sqrt{2\pi}|\beta_2|/(\gamma\tau)$ because $\mathrm{d}C/\mathrm{d}z = 0$ from Eq. (19) for this energy when $\Omega = 0$. In all cases, the solid line shows the

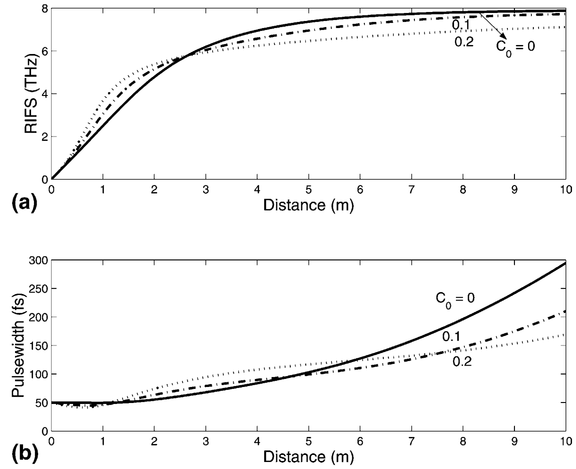

Fig. 3. Evolution of Raman-induced frequency shift (a) and pulse width (b) when Gaussian-shaped pulses with $\lambda = 1.55$ μm and $T_0 = 50$ fs propagate inside a 10-m-long fiber exhibiting anomalous dispersion ($D = 4$ ps/km nm). The input chirp parameter C_0 varies in the range 0–0.2 for the three curves.

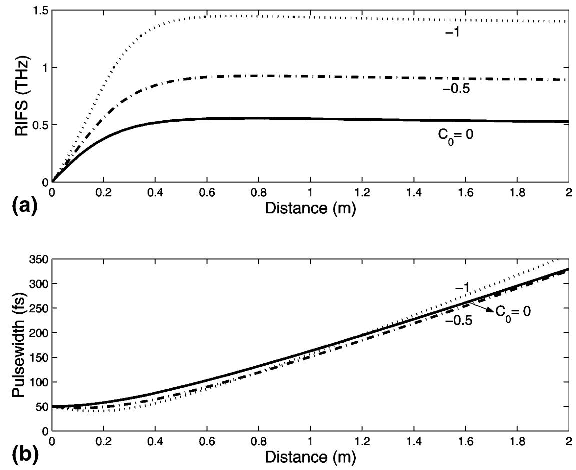

Fig. 4. Same as in Fig. 3 except that fiber exhibits normal dispersion ($D = -4$ ps/km nm). The input chirp parameter C_0 varies in the range −1 to 0 for the three curves.

case when $C_0 = 0$. In the case of anomalous dispersion (Fig. 3), the Gaussian pulse maintains its width up to 2 m, similar to the soliton case, and then broadens because of the RIFS and TOD effects. As expected, Ω increases linearly first and then saturates. Interestingly, the Gaussian pulses acquires a slightly larger RIFS compared with the 'sech' pulses as dispersion-induced broadening depends somewhat on the pulse shape. The dashed and dash-dotted lines show the effects of a positive initial chirp. Since $\beta_2 C_0 < 0$, the pulse undergoes an initial narrowing stage before broadening. A comparison of Figs. 1 and 3 shows that the pulses with nonzero initial chirp experiences the compression stage twice. We attribute this to the imbalance between the dispersive and nonlinear effects in the case of initially chirped pulses. Other qualitative feature are also similar in the two cases. The case of normal dispersion shown in Fig. 4 is quite similar to the results in Fig. 2 obtained for "sech" pulses. For chirp-free Gaussian pulses (solid line), RIFS saturates to a relatively small value of 0.5 THz. However, this value can be increased by applying a negative chirp such that $\beta_2 C_0 < 0$.

5. Conclusions

This paper shows that the RIFS resulting from intrapulse Raman scattering is a general phenomenon that occurs for all pulses both in the normal and anomalous dispersion regimes of an optical fiber. The variational method cannot be used to calculate RIFS because of the dissipative nature of the Raman effect. For this reason, we use the moment method and apply it to the cases of "sech" and Gaussian pulse shapes. The results show that the RIFS depends not only on the width but also on the frequency chirp associated with the optical pulse. The RIFS becomes quite large in the case of ultrashort pulses because, as seen in Eqs. (19) and (26), it depends on the local pulse width as τ^{-3} and varies considerably with the history of pulse width changes. Whenever pulse width remains nearly constant along the fiber, RIFS can accumulate to relatively large values. This is the main reason why RIFS can be quite large for solitons. In the case of standard solitons, our expression for RIFS reduces to that of Gordon [3] as long as the RIFS is much smaller than the spectral width of the pulse. However, we show that even optical solitons do not maintain their width when RIFS becomes comparable to or larger than the spectral width of the pulse. Our analysis remains valid in this regime and shows how RIFS saturates to a constant value because of soliton broadening.

We give numerical examples in both the normal and anomalous dispersion regime using a 10-m long fiber in which femtosecond pulses are launched. Although RIFS is generally smaller for normal dispersion compared with the case of anomalous dispersion, it is large enough to be measurable experimentally. We include the effects of TOD and self-steepening in our analysis and show that even though the TOD does not appear directly in our expression for RIFS, it does affect the RIFS through the frequency chirp. The main limitation of our analysis stems from the assumption that the pulse maintains its shape even though its width may change and it may become chirped. Our results should be used with caution if the pulse shape is known to change significantly during propagation.

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